Probing Large Extra Dimensions Using Top Production in Photon-Photon Collisions

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ABSTRACT

Theories with large extra dimensions predict an infinite tower of Kaluza-Klein states in the 1 TeV range, which can consequently have significant implications for experimental observables. One such observable, which gets affected by the exchange of spin-2 Kaluza-Klein particles, is the $t\bar{t}$ production cross-section in photon-photon collisions at NLC energies. We study this process and obtain bounds on the effective quantum gravity scale M_S between 1600 and 5400 GeV (depending on the centre-of-mass energy). We show that the use of polarisation will further strengthen these bounds.

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Recently there has been tremendous interest in the study of the effects of large compact extra dimensions [1]. In such a scenario, the effects of gravity could become large at very low scales ($\sim$ TeV). Starting from a string theory in 10 dimensions [2, 3], the effective low-energy theory is obtained by compactifying to (3+1) dimensions, in such a way that n of these extra dimensions are compactified to a common scale R which is large, while the remaining dimensions are compactified to scales of the order of the inverse Planck scale. In such a scenario, the Standard Model (SM) particles (corresponding to open strings) live on a 3-brane and are, therefore, confined to the $(3+1)$-dimensional spacetime. On the other hand, the gravitons (corresponding to closed strings) propagate in the $(4+n)$-dimensional bulk. The low-energy scale M_S is related to the Planck scale by [1]

$$M_{\mathrm{P}}^2 = M_S^{n+2} R^n \ , \tag{1}$$

so that we can choose M_S to be of the order of a TeV and thus get around the hierarchy problem. It then follows that $R = 10^{32/n-19}$ m, and so we find that M_S can be arranged to be a TeV for any value $n > 1$. Deviations from classical gravity can become apparent at these surprisingly low values of energy. For example, for $n = 2$ the compactified dimensions are of the order of 1 mm, just below the experimentally tested region for the validity of classical gravity and within the possible reach of ongoing experiments [4]. In fact, it has been shown [5] that it is possible to construct a phenomenologically viable scenario with large extra dimensions, which can survive the existing astrophysical and cosmological constraints. There have also been several interesting studies of the implications of the large Kaluza-Klein dimensions for gauge coupling unification [6]. For some early papers on large Kaluza-Klein dimensions, see Ref. [8, 9] and for recent investigations on different aspects of the TeV scale quantum gravity scenario and related ideas, see Ref. [10].

Below the scale M_S [11, 12, 13], we have an effective theory with an infinite tower of massive Kaluza-Klein states. which contain spin-2, spin-1 and spin-0 excitations. The only states that contribute to low-energy phenomenology in an important manner are the spin-2 Kaluza-Klein states i.e. the infinite tower of massive graviton states in the effective theory. For graviton momenta smaller than the scale M_S, the gravitons couple to the SM fields via a (four-dimensional) induced metric $g_{\mu\nu}$. The interactions of the SM particles with the graviton, $G_{\mu\nu}$, can be derived from the following Lagrangian [13, 12] :

$$\mathcal{L} = -\frac{1}{\bar{M}_P} G_{\mu\nu}^{(j)} T^{\mu\nu} \ , \tag{2}$$

where j labels the Kaluza-Klein mode, $\bar{M}_P = M_P/\sqrt{8\pi}$ and $T^{\mu\nu}$ is the energy-momentum tensor. In spite of being suppressed by $1/\bar{M}_P$, the effects of these couplings are significant for collider energies because of the fact that the sum over the tower of graviton states cancels the dependence on $\bar{M}_P$ and, instead, provides a suppression of the order of M_S.

There have been several studies exploring the consequences of the above effective Lagrangian for experimental observables at high-energy colliders. The gravitons can be directly produced at e^+e^- or hadron colliders leading to spectacular single photon + missing energy or monojet + missing energy signatures [14, 15]. Non-observation of these modes yields bounds which are around 500 GeV to 1.2 TeV at LEP2 [14, 15] and around 600 GeV to 750 GeV at Tevatron (for n between 2 and 6) [14]. These studies have been extended to the Large Hadron Collider (LHC) and to high-energy e^+e^- collisions at

the Next Linear Collider (NLC). Other than looking for the production of gravitons, one can also study the effects of the exchange of virtual gravitons in the intermediate state on experimental observables. Graviton exchange in $e^+e^- \to f\bar{f}$ and in high-mass dilepton production [16, 17], in $t\bar{t}$ production [18] at the Tevatron and the LHC, in deep-inelastic scattering at HERA [19], and in jet production at the Tevatron [20] have been studied. Virtual effects in dilepton production at Tevatron yields a bound of around 950 GeV [16] while a more thorough statistical analysis increases this bound by about 100 GeV [17], $t\bar{t}$ production at Tevatron yields a bound of about 650 GeV [18], while from deep-inelastic scattering a bound of 550 GeV results [19]. Jet production at the Tevatron yields strong bounds of about 1.2 TeV [20]. At LHC, it is expected that $t\bar{t}$ production can be used to explore a range of M_S values upto 4 TeV [18]. More recently, fermion pair production and gauge boson production in e^+e^- collisions at LEP2 and NLC and in $\gamma\gamma$ collisions at the NLC [21, 22, 23, 24, 25] have been studied. Associated production of gravitons with gauge bosons and virtual effects in gauge boson pair production at hadron colliders have also been studied [26]. Diphoton signals and global lepton-quark neutral current constraints have also been studied [27]. There have also been papers discussing the implications of the large dimensions for higgs production [28, 29] and electroweak precision observables [30]. Astrophysical constraints, like bounds from energy loss for supernovae cores, have also been discussed [31].

The present work concentrates on studying the effects of large extra dimensions in top production in photon-photon collisions at the NLC. The photons are produced in the Compton back-scattering of a highly monochromatic low-energy laser beam off a high energy electron beam [32]. Control over the e^- and laser beam parameters allow for control over the parameters of the photon-photon subprocess. The physics potential of the NLC in the photon-photon mode is manifold and these experiments are planned over several steps of energy spanning the range between 500 GeV and 1.5 TeV. These experiments also provide a great degree of precision because of the relatively clean initial state, and indeed the degree of precision can be enhanced by using polarised initial electrons and laser beams. The precision that is possible in these experiments and the high energies that are planned to be accessed provide an ideal testing ground of the SM and a very effective probe of possible physics that may lie beyond the SM. In particular, the physics of large extra dimensions can be studied in $\gamma\gamma$ collisions at the NLC with the exciting possibility of discovering this source of new physics, or (in the absence of any evidence for its existence) to put stringent bounds on the parameters of this model. We study the effects of virtual graviton exchange on the $t\bar{t}$ production cross-section in $\gamma\gamma$ collisions with the above physics goal in mind.

The simplest way to approach the photon-photon scattering process in these experiments is to think of it as analogous to the parton model so that the basic scattering is described by a $\gamma\gamma$ scattering subprocess, with each γ resulting from the electron-laser back scattering. The energy of the back-scattered photon, E_γ, follows a distribution characteristic of the Compton scattering process and can be written in terms of the dimensionless ratio $x = E_\gamma/E_e$. It turns out that the maximum value of x is about 0.82 so that provides the upper limit on the energy accessible in the $\gamma\gamma$ sub-process. The subprocess cross-section is convoluted with the luminosity functions, $f_\gamma^i(x)$, which provide information on the photon flux produced in Compton scattering of the electron and laser beams [33].

For the case of $t\bar{t}$ production that we consider in this paper, the above discussion implies that we need only calculate the cross-section for the $\gamma\gamma \to t\bar{t}$ process and then convolute this with the

luminosity functions to get the full cross-section. The cross-section for the $\gamma\gamma \to t\bar{t}$ process has the usual t- and u-channel SM contributions, but in addition, we also have the s-channel exchange of virtual spin-2 Kaluza-Klein particles. In the following, we refer to this contribution as the non-SM (NSM) contribution. We calculate the cross-section for the polarised case and then obtain the unpolarised cross-section by summing over the polarisations of the initial photons. The polarisation of each of the photon is a function of the polarisations of the initial electron and laser beams and it is only the latter that can be fixed in the experiment. When we present our results for the polarised case, we will do so for a fixed choice of these initial polarisations.

For the polarised case, the cross-section takes on the following factorised form

$$d\sigma(\lambda_{e1}, \lambda_{e2}, \lambda_{\ell1}, \lambda_{\ell2}) = \int dx_1 \int dx_2 \ f^1_\gamma(x_1, \xi) \ f^2_\gamma(x_2, \bar{\xi})$$
$$\times \left(d\hat{\sigma}_{00} + \xi_2 \ \bar{\xi}_2 \ d\hat{\sigma}_{22} + \xi_2 \ d\hat{\sigma}_{02} + \bar{\xi}_2 \ d\hat{\sigma}_{20} \right) \ d\Gamma, \tag{3}$$

where $d\Gamma$ includes the phase space and flux factor. ξ and $\bar{\xi}$ are the Stokes parameters of the first and second photons respectively. For example, $f^1_\gamma(x_1, \xi)$ is the distribution of photons for a given polarisation $\lambda_{\ell1}$ of the laser beam and a polarisation λ_{e1} of the electron. $d\hat{\sigma}_{ij}$ is related to the polarised $\gamma \gamma \to t \bar{t}$ matrix elements as follows

$$d\hat{\sigma}_{00} = \frac{1}{4} \sum_{\lambda_1 \lambda_2} |M(\lambda_1, \lambda_2)|^2,$$

$$d\hat{\sigma}_{22} = \frac{1}{4} \sum_{\lambda_1 \lambda_2} \lambda_1 \lambda_2 |M(\lambda_1, \lambda_2)|^2,$$

$$d\hat{\sigma}_{20} = \frac{1}{4} \sum_{\lambda_1 \lambda_2} \lambda_1 |M(\lambda_1, \lambda_2)|^2,$$

$$d\hat{\sigma}_{02} = \frac{1}{4} \sum_{\lambda_1 \lambda_2} \lambda_2 |M(\lambda_1, \lambda_2)|^2,$$

where $\lambda_{1,2}$ denote the circular polarisation of the photons 1,2 respectively.

Summing over the polarisations of the electron and laser beam the only term that survives in the above expression is $d\sigma_{00}$. The unpolarised differential cross section can be written as

$$\frac{d^3\sigma}{dp_T^2 d \ y} = \int dx_1 \int dx_2 \ f_\gamma(x_1) \ f_\gamma(x_2) \ \hat{s}\frac{d\hat{\sigma}}{d\hat{t}} \ \delta(\hat{s} + \hat{t} + \hat{u} - 2m_t^2), \tag{4}$$

where the hatted variables corresponds to the $\gamma \gamma \to t \bar{t}$ subprocess and m_t is the mass of the top quark.

$$\frac{d\hat{\sigma}}{d\hat{t}} = \left(\frac{d\hat{\sigma}}{d\hat{t}}\right)_{SM} + \left(\frac{d\hat{\sigma}}{d\hat{t}}\right)_{NSM}, \tag{5}$$

3

$$\left(\frac{d\hat{\sigma}}{d\hat{t}}\right)_{NSM}(\lambda_1, \lambda_2) = -\frac{3}{2\hat{s}^2}[1 - \lambda_1\lambda_2]\left(\frac{\pi\lambda^2}{M_s^8} - \frac{4\pi\alpha e_q^2\lambda}{M_s^4}\frac{1}{(m_t^2 - \hat{t})(m_t^2 - \hat{u})}\right)$$
$$\times\left\{2m_t^8 - 8m_t^6\hat{t} - 2m_t^2\hat{t}(\hat{s} + 2\hat{t})^2 + m_t^4(\hat{s}^4 + 4\hat{s}\hat{t} + 12\hat{t}^2)\right.$$
$$\left.+\hat{t}(\hat{s}^3 + 3\hat{s}^2\hat{t} + 4\hat{s}\hat{t}^2 + 2\hat{t}^3)\right\}. \tag{6}$$

Before we discuss the analysis that we have done, we would like to point out that the non-SM contribution involves two new parameters: the effective string scale M_S and λ which is the effective coupling at M_S. λ is expected to be of $\mathcal{O}(1)$, but its sign is not known *a priori*. In our work we will explore the sensitivity of our results to the choice of the sign of λ.
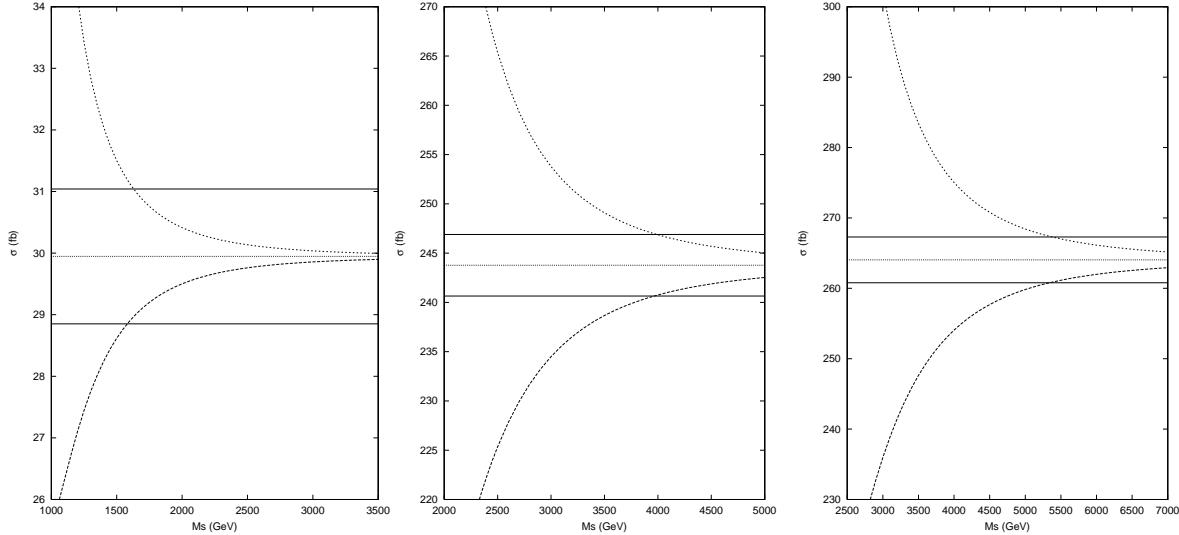

Fig.1: $t\bar{t}$ production cross-section (in fb) as a function of M_S for $\sqrt{s}$ = 500 GeV (a), 1000 GeV (b) and 1500 GeV (c). The upper (lower) curves are for $\lambda = -1(+1)$. The lines show the SM value and the 2σ upper and lower limits.

We begin with the results for the unpolarised case. In Fig. 1, we have plotted the unpolarised (integrated) cross-section as a function of M_S for three different values of the initial e^+e^- C.M. energy i.e. $\sqrt{s}$ = 500, 1000, 1500 GeV (shown in Fig. 1 (a), (b) and (c) respectively). Also shown in the figure are the SM value of the cross-section and the 2σ upper and lower limits. For obtaining the latter, we have assumed purely statistical errors and assuming an integrated luminosity of 100 fb^{-1}. We find, first of all, that the curves are symmetrical with respect to λ. The 2σ limits that we obtain for $\sqrt{s}$ = 500, 1000, 1500 GeV are 1600, 4000 and 5400 GeV, respectively. At the higher C.M. energies of $\sqrt{s}$ = 1000, 1500 GeV that we have considered, the $t\bar{t}$ production cross-section is large enough and with

the integrated luminosity that we have assumed a large number of $t\bar{t}$ events results. The integrated cross-section itself becomes a good discriminator of the new physics effects.

In Fig. 2 we have plotted the rapidity distribution for $\sqrt{s} = 1000$ GeV, in order to consider whether the use of differential quantities will further enhance the sensitivity of the process under consideration to the effects of the new physics. We find that the difference between the SM and the SM+NSM distributions to be quite significant, especially when we concentrate in the central regions of rapidity. Though we refrain from making a quantitative estimate of the bound that would result (for such an estimate would be premature without knowing further experimental details), it is clear from Fig. 2 that the rapidity distribution can be used to improve the bound that would result from looking only at the integrated cross-section.
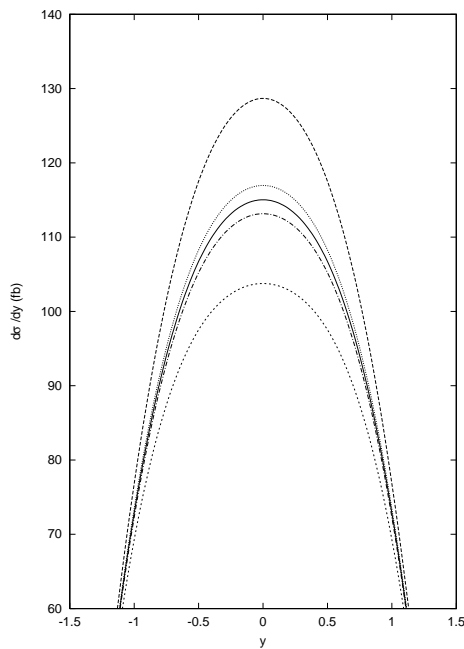

Fig.2: The rapidity distribution for $\sqrt{s} = 1000$ GeV. The solid curve is the SM distribution. The two curves below the SM curve are for $M_S = 1.2$, 1.8 TeV respectively and for $\lambda = 1$, while the curves above the SM curve for $\lambda = -1$.

For the polarised case, we study the effects of the NSM contribution for different choices of the initial electron and laser beams. For a given choice of the e^- and laser polarisation, the photon polarisation is fixed once the x value is known. The latter polarisation is therefore dependent crucially on the luminosity functions and it is only on the polarisation of the electron and the laser beams that we have a direct handle. The efficacy of polarisation as a discriminator of the new physics is, however, apparent more at the level of the $\gamma\gamma$ sub-process. As we scan over the different choices of initial beam polarisations, we find that for certain choices there is hardly any sensitivity to the new physics

$(\lambda_{e1}\ \lambda_{e2}\ \lambda_{\ell 1}\ \lambda_{\ell 2})$	$\sqrt{s}$ (GeV)	M_s (GeV)
	500	1950
$(+\quad -\quad -\quad -\ )$	1000	4600
	1500	6000
	500	2500
$(+\quad -\quad -\quad +)$	1000	4800
	1500	6400

Table 1: 2σ limits on Ms (GeV) with polarized initial beams $\lambda_{e1}\ \lambda_{e2}\ \lambda_{\ell 1}\ \lambda_{\ell 2}$ are the helicities of initial electron and laser beams and $\sqrt{s}$ is the e^-e^- centre of mass energy. An e^-e^- Luminosity of $100 fb^{-1}$ is assumed.

i.e. the SM and the SM+NSM cross-sections are more or less the same. However, large differences are realised for certain other choices *viz.* for the cases $(+,-,-,-)$ and $(+,-,-,+)$, where these represent the polarisations $(\lambda_{e1}, \lambda_{e2}, \lambda_{\ell 1}, \lambda_{\ell 2})$. These polarization combinations are very sensitive to the NSM contribution while other combinations are not so. This has to do with the fact that the NSM contribution (see eqn.6) is significant when λ_1 and λ_2 are with opposite signs. For hard photons, the sign of λ is the same as that of the helicity of the initial electron. Therefore the combination that has $\lambda_1\lambda_2$ −ve is the one with electron beams with opposite helicities ($\lambda_{e1} = -\lambda_{e2}$). The luminosity function is such that, for the combination with $\lambda_e\lambda_l = -1$, it peaks in the high energy region. Therefore, the sensitivity of the combination with both the beams having $\lambda_{e1}\lambda_{l1} = -1 = \lambda_{e2}\lambda_{l2}$ is the best while the combination with $\lambda_{e1}\lambda_{l1} = 1 = \lambda_{e2}\lambda_{l2}$ is least sensitive. The other combination has a sensitivity which is intermediate.

In Table 1, we show the M_S limits for these polarisation choices for the three different C.M. energies that we have considered. As can be clearly seen, the use of polarisation enhances the bounds on M_S quite significantly by several 100 GeV in each case.

In summary, the effects of virtual spin-2 particle exchange in scenarios with large extra Kaluza-Klein dimensions can be significantly tested in $t\bar{t}$ production in $\gamma\gamma$ collisions at the NLC. The unpolarised integrated rate in itself turns out to be a sensitive discriminator for this scenario and the bounds on the effective low-energy quantum gravity scale that result from the analysis of the integrated rate is between approximately 1600 GeV and 5400 GeV depending upon the initial e^+e^- C.M. energy. These bounds can be strengthened by studying rapidity distributions or by tuning the polarisations of the initial electron and laser beams.

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